

Quark-flavour violating Higgs decays to charm and bottom pairs in the MSSM

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Abstract

We calculate the decay width of $h^0 \rightarrow b\bar{b}$ in the Minimal Supersymmetric Standard Model (MSSM) with quark-flavour violation (QFV) at full one-loop level. The effect of $\tilde{c}-\tilde{t}$ mixing and $\tilde{s}-\tilde{b}$ mixing is studied taking into account the constraints from the B-meson data. We discuss and compare in detail the decays $h^0 \rightarrow c\bar{c}$ and $h^0 \rightarrow b\bar{b}$ within the framework of the perturbative mass insertion technique using the Flavour Expansion Theorem. The deviation of both decay widths from the Standard Model results can be quite large. While in $h^0 \rightarrow c\bar{c}$ it is almost entirely due to the flavour violating part of the MSSM, in $h^0 \rightarrow b\bar{b}$ it is mainly due to the flavour conserving part. Nevertheless, $\Gamma(h^0 \rightarrow b\bar{b})$ can fluctuate up to $\sim 7\%$ due to QFV chargino exchange with large $\tilde{c}-\tilde{t}$ mixing.

Keywords: High Energy Physics, Supersymmetry, Higgs Physics, Flavour Physics

1. Introduction

So far, the Higgs boson properties measured at the LHC experiments are consistent with the Standard Model (SM) predictions. Deviations are, however, not yet excluded and could indicate physics beyond the SM. For instance, in the Minimal Supersymmetric Standard Model (MSSM), the discovered Higgs boson can well be the lightest neutral Higgs boson, h^0 , with a mass of 125 GeV and SM-like couplings. Non-minimal quark-flavour violation (QFV) in the squark sector of the MSSM can additionally affect the Higgs interactions at one-loop level, still being consistent with the B-physics constraints. We study two important decays of the Higgs boson: into a pair of bottom quarks and into a pair of charm quarks at one-loop level in the MSSM with general squark mixing. We consider mainly mixing between the two heavy generations up- and down-squarks, i.e. $\tilde{c}_{L,R}-\tilde{t}_{L,R}$ and $\tilde{s}_{L,R}-\tilde{b}_{L,R}$ mixing. We investigate numerically the influence of such mixing on the Higgs properties, taking into account the constraints on quark-flavour mixing from B-physics. This talk is based on Ref. [1], for more details see the original publication.

2. QFV in the squark sector of the MSSM

We define the QFV parameters in the up-type squark sector of the MSSM as follows:

$$\delta_{\alpha\beta}^{LL} \equiv M_{Q\alpha\beta}^2 / \sqrt{M_{Q\alpha\alpha}^2 M_{Q\beta\beta}^2}, \quad (1)$$

$$\delta_{\alpha\beta}^{uRR} \equiv M_{U\alpha\beta}^2 / \sqrt{M_{U\alpha\alpha}^2 M_{U\beta\beta}^2}, \quad (2)$$

$$\delta_{\alpha\beta}^{uRL} \equiv (v_2 / \sqrt{2}) T_{U\alpha\beta} / \sqrt{M_{U\alpha\alpha}^2 M_{Q\beta\beta}^2}, \quad (3)$$

where $\alpha, \beta = 1, 2, 3$ ($\alpha \neq \beta$) denote the quark flavours u, c, t , and $v_2 = \sqrt{2} \langle H_2^0 \rangle$. $M_{Q,U}$ are the hermitian soft SUSY-breaking squark mass matrices and T_U are the soft SUSY-breaking trilinear coupling matrices of the up-type squarks. These parameters enter the left-left, right-right and left-right blocks of the 6×6 up-type squark mass matrix in the super-CKM basis [2],

$$\mathcal{M}_{\tilde{u}}^2 = \begin{pmatrix} \mathcal{M}_{\tilde{u},LL}^2 & \mathcal{M}_{\tilde{u},LR}^2 \\ \mathcal{M}_{\tilde{u},RL}^2 & \mathcal{M}_{\tilde{u},RR}^2 \end{pmatrix}. \quad (4)$$

The different blocks in eq. (4) are given by

$$\mathcal{M}_{\tilde{u},LL}^2 = V_{\text{CKM}} M_Q^2 V_{\text{CKM}}^\dagger + D_{\tilde{u},LL} \mathbf{1} + \hat{m}_u^2,$$

$$\begin{aligned}\mathcal{M}_{\tilde{u},RR}^2 &= M_U^2 + D_{\tilde{u},RR} \mathbf{1} + \hat{m}_u^2, \\ \mathcal{M}_{\tilde{u},RL}^2 &= \mathcal{M}_{\tilde{u},LR}^{2\dagger} = \frac{v_2}{\sqrt{2}} T_U - \mu^* \hat{m}_u \cot \beta,\end{aligned}\quad (5)$$

where μ is the higgsino mass parameter, $\tan \beta$ is the ratio of the vacuum expectation values of the neutral Higgs fields v_2/v_1 , with $v_{1,2} = \sqrt{2} \langle H_{1,2}^0 \rangle$, and \hat{m}_u is the diagonal mass matrix of the up-type quarks. Furthermore, $D_{\tilde{u},LL} = \cos 2\beta m_Z^2 (T_3^u - e_u \sin^2 \theta_W)$ and $D_{\tilde{u},RR} = e_u \sin^2 \theta_W \times \cos 2\beta m_Z^2$, with T_3^u and e_u being the isospin and electric charge of the up-type quarks (squarks), respectively, and θ_W is the weak mixing angle. V_{CKM} is the Cabibbo-Kobayashi-Maskawa matrix, which we approximate with the unitary matrix. The up-squark mass matrix is diagonalized by the 6×6 matrices $U^{\tilde{u}}$, such that

$$U^{\tilde{u}} \mathcal{M}_{\tilde{u}}^2 (U^{\tilde{u}})^\dagger = \text{diag}(m_{\tilde{u}_1}^2, \dots, m_{\tilde{u}_6}^2), \quad (6)$$

with the mass hierarchy $m_{\tilde{u}_1} < \dots < m_{\tilde{u}_6}$. The physical mass eigenstates $\tilde{u}_i, i = 1, \dots, 6$ are given by $\tilde{u}_i = U_{i\alpha}^{\tilde{u}} \tilde{u}_{0\alpha}$. The QFV parameters of the down squark sector are defined analogously, see Ref. [1].

We mainly focus on the $\tilde{c}_R - \tilde{t}_L, \tilde{c}_L - \tilde{t}_R, \tilde{c}_R - \tilde{t}_R$, and $\tilde{c}_L - \tilde{t}_L$ mixing, which is described by the QFV parameters $\delta_{23}^{uRL}, \delta_{23}^{uLR} \equiv (\delta_{32}^{uRL})^*, \delta_{23}^{uRR}$, and δ_{23}^{LL} , respectively. The $\tilde{t}_R - \tilde{t}_L$ mixing is described by the quark-flavour conserving (QFC) parameter δ_{33}^{uRL} . All the QFV and QFC parameters are assumed to be real.

3. The processes

The decay width of $h^0 \rightarrow q\bar{q}$, with $q = c, b$, including one-loop contributions, can be written as

$$\Gamma(h^0 \rightarrow q\bar{q}) = \Gamma^{\text{tree}}(h^0 \rightarrow q\bar{q}) + \delta\Gamma^{\text{1loop}}(h^0 \rightarrow q\bar{q}) \quad (7)$$

with the tree-level decay width

$$\Gamma^{\text{tree}}(h^0 \rightarrow q\bar{q}) = \frac{N_C}{8\pi} m_{h^0} (s_1^q)^2 \left(1 - \frac{4m_q^2}{m_{h^0}^2}\right)^{3/2}, \quad (8)$$

where $N_C = 3$, m_{h^0} is the on-shell mass of h^0 and the tree-level coupling s_1^q is given by

$$s_1^{c,b} = \mp g \frac{m_{c,b}}{2m_W} \frac{\cos \alpha}{\sin \beta} \left(\frac{\sin \alpha}{\cos \beta} \right) = \mp \frac{h_{c,b}}{\sqrt{2}} \cos \alpha (\sin \alpha), \quad (9)$$

α is the mixing angle of the two CP-even Higgs bosons, h^0 and H^0 .

3.1. Gluino contribution to $h^0 \rightarrow c\bar{c}$

For the calculation of $\delta\Gamma^{\text{1loop}}(h^0 \rightarrow b\bar{b})$ we proceed in a way analogous to the calculation of $\delta\Gamma^{\text{1loop}}(h^0 \rightarrow c\bar{c})$ in Ref. [3]. The corresponding loop diagrams are shown in Fig. 2 of [3], with the replacements: $c \leftrightarrow b$ and $\tilde{u} \leftrightarrow \tilde{d}$. The dominant supersymmetric (SUSY) contribution is due to gluino and chargino exchange, which also contribute to the self-energy of the b-quark.

As in Ref. [3], in our calculation we employ the $\overline{\text{DR}}$ renormalisation scheme, with the Lagrangian input parameters, defined at the scale $Q = 1$ TeV. The shifts from the $\overline{\text{DR}}$ masses and fields to the physical scale-independent quantities are obtained using on-shell renormalisation conditions. To assure infrared (IR) convergence we include the real gluon/photon radiation contributions as well.

Furthermore, we compare and recalculate in the mass insertion (MI) technique $\Gamma(h^0 \rightarrow b\bar{b})$ and $\Gamma(h^0 \rightarrow c\bar{c})$, as previously studied in Ref. [3]. In this context we often refer to the one-loop representation

$$\Gamma(h^0 \rightarrow b\bar{b}) = \Gamma^{\text{g,impr}} + \delta\Gamma^{\text{g}} + \delta\Gamma^{\text{EW}}, \quad (10)$$

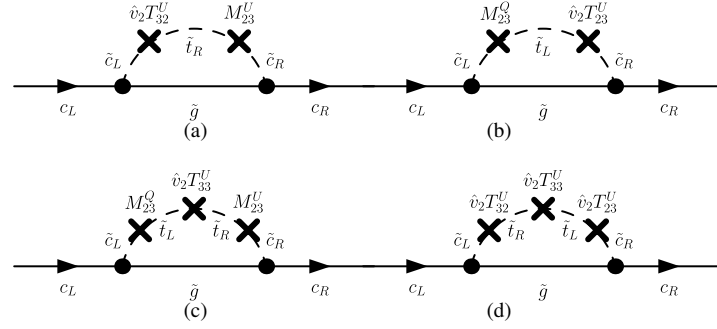
where $\Gamma^{\text{g,impr}}$ includes the tree-level and the gluon one-loop contribution (see eq.(55) in [3]), $\delta\Gamma^{\text{g}}$ is the gluino one-loop contribution, and $\delta\Gamma^{\text{EW}}$ is the electroweak one-loop contribution.

4. Mass insertion technique

The perturbative interaction between the Higgs and the squarks is explicitly proportional to the soft SUSY-breaking trilinear coupling matrices, $T_{U,D}$. However, the dependence on the soft SUSY-breaking mass matrices, $\mathcal{M}_{\tilde{u},\tilde{d}}^2$ is hidden in the squark mixing matrices, $U^{\tilde{u},\tilde{d}}$, which makes the analysis complicated. An effective approach using the mass insertion (MI) approximation gives access to the explicit dependences on these QFV parameters and allows an analytic approach to study the QFV effects. In our calculations we exploit the Flavour Expansion Theorem (FET) [4].

In the following we briefly review the main suggestion of the MI approximation. If $X = U_{iA}^q U_{iB}^{q*} B_0(0, m^2, m_{q_i}^2)$, where B_0 is the two-point function given in terms of mass eigenstates, and U^q are the rotation matrices defined with eq. (6) ($A \neq B$), then X can be expanded into mass insertions (MIs) by the FET [4]:

$$\begin{aligned}X &= M_{AB}^I b_0(1, m^2, \{M_{AA}, M_{BB}\}) \\ &+ M_{Ai}^I M_{iB}^I b_0(2, m^2, \{M_{AA}, M_{ii}, M_{BB}\})\end{aligned}$$

Figure 1: Quark-flavour violating mass insertions to the charm quark self-energy with gluino, corresponding to T_2 and T_3 in eq. (17).

$$\begin{aligned}
 & + M_{Ai}^I M_{ij}^I M_{jB}^I b_0(3, m^2, \{M_{AA}, M_{ii}, M_{jj}, M_{BB}\}) \\
 & + M_{Ai}^I M_{ij}^I M_{jk}^I M_{kB}^I b_0(4, m^2, \\
 & \quad \{M_{AA}, M_{ii}, M_{jj}, M_{kk}, M_{BB}\}) + \dots, \quad (11)
 \end{aligned}$$

The insertions are represented by the elements of the matrix M^I , with $M_{ii}^I = 0$. The generalized b_0 functions used in eq. (11), where the first argument shows the number of insertions, can be written recursively as [4]

$$\begin{aligned}
 b_0(1, a, \{b, c\}) &= \frac{b_0(a, b) - b_0(a, c)}{b - c}, \\
 b_0(2, a, \{b, c, d\}) &= \frac{b_0(1, a, \{b, c\}) - b_0(1, a, \{b, d\})}{c - d}, \\
 b_0(3, a, \{b, c, d, e\}) &= \frac{b_0(2, a, \{b, c, d\}) - b_0(2, a, \{b, c, e\})}{d - e}, \\
 b_0(4, a, \{b, c, d, e, f\}) &= \frac{b_0(3, a, \{b, c, d, e\}) - b_0(3, a, \{b, c, d, f\})}{e - f}, \quad (12)
 \end{aligned}$$

where

$$\begin{aligned}
 b_0(a, b) &\equiv B_0(0, a, b) = \\
 &= \frac{b \log\left(\frac{b}{Q^2}\right) - a \log\left(\frac{a}{Q^2}\right)}{a - b} + \Delta + 1, \quad (13)
 \end{aligned}$$

with the renormalisation scale Q and the UV-divergence parameter Δ .

4.1. Gluino contribution to $h^0 \rightarrow c\bar{c}$

In order to demonstrate how the mass insertion approximation works we calculate the self-energy of the c -quark Σ_c with \tilde{g} and \tilde{u}_i in the loop. The relevant LR -part $\Sigma_c^{LR} = \Sigma_c^{RL}$ reads

$$\Sigma_c^{LR, \tilde{g}} = -\frac{2\alpha_s}{3\pi} \frac{m_{\tilde{g}}}{m_c} \sum_{i=1}^6 U_{i2}^{\tilde{u}*} U_{i5}^{\tilde{u}} B_0(m_c^2, m_{\tilde{g}}^2, m_{\tilde{u}_i}^2), \quad (14)$$

where we assume that the squared \tilde{u} -mass matrix $\mathcal{M}_{\tilde{u}}^2 \equiv M_{ij}$ (see eq. (4)) is in the form

$$\mathcal{M}_{\tilde{u}}^2 = \begin{pmatrix} M_{11}^{LL} & 0 & 0 & 0 & 0 & 0 \\ 0 & M_{22}^{LL} & M_{23}^Q & 0 & 0 & \hat{v}_2 T_{32}^U \\ 0 & M_{23}^Q & M_{33}^{LL} & 0 & \hat{v}_2 T_{23}^U & \hat{v}_2 T_{33}^U \\ 0 & 0 & 0 & M_{11}^{RR} & 0 & 0 \\ 0 & 0 & \hat{v}_2 T_{23}^U & 0 & M_{22}^{RR} & M_{23}^U \\ 0 & \hat{v}_2 T_{32}^U & \hat{v}_2 T_{33}^U & 0 & M_{23}^U & M_{33}^{RR} \end{pmatrix}, \quad (15)$$

with $\hat{v}_2 = v \sin \beta / \sqrt{2} \sim 170$ GeV, and the QFV elements of the 3×3 matrices $M_{\tilde{Q}}^2$ and $M_{\tilde{U}}^2$ are denoted by M_{ij}^Q and M_{ij}^U , respectively. Using the FET (eq. (11)) we get

$$m_c \Sigma_c^{LR, \tilde{g}} = -\frac{2\alpha_s}{3\pi} m_{\tilde{g}} (T_2 + T_3 + T_4 + \dots), \quad (16)$$

where the QFV contributions read

$$\begin{aligned}
 T_2 &= \hat{v}_2 T_{32}^U M_{23}^U b_0(2, m_{\tilde{g}}^2, \{M_{22}^{LL}, M_{33}^{RR}, M_{22}^{RR}\}) \\
 &+ \hat{v}_2 T_{23}^U M_{23}^Q b_0(2, m_{\tilde{g}}^2, \{M_{22}^{LL}, M_{33}^{LL}, M_{22}^{RR}\}) \\
 T_3 &= \hat{v}_2 T_{33}^U (M_{23}^Q M_{23}^U + 3\rho \hat{v}_2^2 T_{23}^U T_{32}^U) \\
 &\times b_0(3, m_{\tilde{g}}^2, \{M_{22}^{LL}, M_{33}^{LL}, M_{33}^{RR}, M_{22}^{RR}\}), \quad (17)
 \end{aligned}$$

with $\rho = 0$. The corresponding to the terms T_2 and T_3 graphs are shown in Figs. 1(a) and 1(b) or 1(c) and 1(d), respectively. Note that there is no contribution with no mass insertion because of the helicity flip, and also practically no contribution with only one insertion, because $T_{22}^U \approx 0$. Thus, all terms in eq. (17) are quark-flavour violating.

The vertex contributions with \tilde{g} , \tilde{u}_i^* and \tilde{u}_j in the loop, defined by $\mathcal{L} = -h^0 \bar{c} (c_L^\nu P_L + c_R^\nu P_R) c$, are calculated in an analogous way. Here we only show the result for the mass insertion expansions for the coefficients c_L^ν and c_R^ν , for details see Ref.[1]. In case of real input parameters, $c_L^\nu = c_R^\nu = c^\nu$, we obtain

$$c^\nu = -\frac{2\alpha_s}{3\pi} m_{\tilde{g}} \frac{\cos \alpha}{\sqrt{2}} (T_1^\nu + T_2^\nu + \dots), \quad (18)$$

where

$$\begin{aligned} T_1^v &= T_{32}^U M_{23}^U b_0(2, m_{\tilde{g}}^2, \{M_{22}^{LL}, M_{33}^{RR}, M_{22}^{RR}\}) \\ &+ T_{23}^U M_{23}^Q b_0(2, m_{\tilde{g}}^2, \{M_{22}^{LL}, M_{33}^{LL}, M_{22}^{RR}\}), \\ T_2^v &= T_{33}^U (M_{23}^Q M_{23}^U + 3 \hat{v}_2^2 T_{23}^U T_{32}^U) \\ &\times b_0(3, m_{\tilde{g}}^2, \{M_{22}^{LL}, M_{33}^{LL}, M_{33}^{RR}, M_{22}^{RR}\}). \end{aligned} \quad (19)$$

Comparing the results for the charm self-energy, eqs. (16), (17), and the vertex contribution to $h^0 \rightarrow c\bar{c}$, eqs. (18), (19), we see that $T_2 = T_1^v \hat{v}_2$. The same holds for the term proportional to $T_{33}^U M_{23}^Q M_{23}^U$ in T_3 and T_2^v . Concerning the term proportional to $T_{33}^U T_{23}^U T_{32}^U$, we have a factor 3 in the term T_2^v compared to that in T_3 . Thus we can deduce the result T_3^v from the term T_4 in eq. (17) by adding a prefactor of 3 for all the terms with three T^U elements.

4.2. Gluino and chargino contributions to $h^0 \rightarrow b\bar{b}$

Assuming the squared \tilde{d} -mass matrix $\mathcal{M}_{\tilde{d}}^2 \equiv M_{ij}$ in the form

$$\mathcal{M}_{\tilde{d}}^2 = \begin{pmatrix} M_{11}^{LL} & 0 & 0 & 0 & 0 & 0 \\ 0 & M_{22}^{LL} & M_{23}^Q & 0 & 0 & \hat{v}_1 T_{32}^D \\ 0 & M_{23}^Q & M_{33}^{LL} & 0 & \hat{v}_1 T_{23}^D & M_{33}^{RL} \\ 0 & 0 & 0 & M_{11}^{RR} & 0 & 0 \\ 0 & 0 & \hat{v}_1 T_{23}^D & 0 & M_{22}^{RR} & M_{23}^D \\ 0 & \hat{v}_1 T_{32}^D & M_{33}^{RL} & 0 & M_{23}^D & M_{33}^{RR} \end{pmatrix}, \quad (20)$$

with $M_{33}^{RL} \sim -\mu m_b \tan\beta$, $\hat{v}_1 = v \cos\beta / \sqrt{2}$, and the QFV elements of the 3×3 matrices $M_{\tilde{Q}}^2$ and $M_{\tilde{D}}^2$ denoted by M_{ij}^Q and M_{ij}^D , respectively, for the LR -part of the gluino contribution to the bottom self energy $\Sigma_b^{LR, \tilde{g}}$, defined by the Lagrangian $\mathcal{L} = -\bar{b} \Sigma_b b$, we obtain

$$m_b \Sigma_b^{LR, \tilde{g}} = -\frac{2\alpha_s}{3\pi} m_{\tilde{g}} (T_1^{FC} + T_2^{FV} + T_3^{FC} + T_3^{FV} + \dots) \quad (21)$$

where the quark flavour conserving (FC) and quark flavour violating (FV) contributions read

$$\begin{aligned} T_1^{FC} &= M_{33}^{RL} b_0(1, m_{\tilde{g}}^2, \{M_{33}^{RR}, M_{33}^{LL}\}) \\ T_2^{FV} &= \hat{v}_1 T_{32}^D M_{23}^Q b_0(2, m_{\tilde{g}}^2, \{M_{33}^{RR}, M_{22}^{LL}, M_{33}^{LL}\}) \\ &+ \hat{v}_1 T_{23}^D M_{23}^D b_0(2, m_{\tilde{g}}^2, \{M_{33}^{RR}, M_{22}^{RR}, M_{33}^{LL}\}) \\ T_3^{FC} &= 3^\rho (M_{33}^{RL})^3 \\ &\times b_0(3, m_{\tilde{g}}^2, \{M_{33}^{RR}, M_{33}^{LL}, M_{33}^{RR}, M_{33}^{LL}\}) \\ T_3^{FV} &= (M_{23}^Q)^2 M_{33}^{RL} \\ &\times b_0(3, m_{\tilde{g}}^2, \{M_{33}^{RR}, M_{33}^{LL}, M_{22}^{LL}, M_{33}^{LL}\}) \\ &+ (M_{23}^D)^2 M_{33}^{RL} \\ &\times b_0(3, m_{\tilde{g}}^2, \{M_{33}^{RR}, M_{33}^{RR}, M_{22}^{RR}, M_{33}^{LL}\}) \end{aligned}$$

$$\begin{aligned} &+ 3^\rho (\hat{v}_1)^2 (T_{32}^D)^2 M_{33}^{RL} \\ &\times b_0(3, m_{\tilde{g}}^2, \{M_{33}^{RR}, M_{22}^{LL}, M_{33}^{RR}, M_{33}^{LL}\}) \\ &+ 3^\rho (\hat{v}_1)^2 (T_{23}^D)^2 M_{33}^{RL} \\ &\times b_0(3, m_{\tilde{g}}^2, \{M_{33}^{RR}, M_{33}^{LL}, M_{22}^{RR}, M_{33}^{LL}\}) \end{aligned} \quad (22)$$

with $\rho = 0$. As in the $c\bar{c}$ case, the vertex contribution can be directly deduced from the self energy, with $T_i^{v,x} = \frac{T_i^x}{\hat{v}_1}$, $x = FC, FV$, an additional factor 3 for some terms in T_3^x , $(M_{33}^{RL})^3 \rightarrow 3(M_{33}^{RL})^3$, $(T_{23}^D)^2 M_{33}^{RL} \rightarrow 3(T_{23}^D)^2 M_{33}^{RL}$, $(T_{32}^D)^2 M_{33}^{RL} \rightarrow 3(T_{32}^D)^2 M_{33}^{RL}$, and accordingly $\rho = 1$.

The relevant term for the self-energy calculation of the bottom-quark and for the vertex amplitude with a chargino in the loop is proportional to $c_L^* c_R$, with $c_L = h_b U_{m2}^* U_{i3}^{\tilde{u}*}$ and $c_R = -g V_{m1} U_{i3}^{\tilde{u}*} + h_t V_{m2} U_{i6}^{\tilde{u}*}$, where U and V diagonalize the chargino mass matrix X : $U^* X V^{-1} = M_D = \text{diag}(m_{\tilde{\chi}_1^+}, m_{\tilde{\chi}_2^+})$. Neglecting the term proportional to g and m_b in the loop integrals, we obtain

$$\begin{aligned} m_b \Sigma_b^{LR, \tilde{\chi}^+} &= \frac{h_b h_t}{16\pi^2} \sum_{m=1}^2 \sum_{i=1}^6 m_{\tilde{\chi}_m^+} U_{m2} V_{m2} U_{i3}^{\tilde{u}} U_{i6}^{\tilde{u}*} \\ &\times b_0(m_{\tilde{\chi}_m^+}^2, m_{\tilde{u}_i}^2). \end{aligned} \quad (23)$$

The mass insertions in the \tilde{u}_i line can be deduced from the results for the bottom self-energy with gluino in the loop. Moreover, we can also apply the MI technique to the chargino part ($\sim U_{m2} V_{m2}$) in eq. (23) using linear approximation. Finally, we obtain the approximate result

$$m_b \Sigma_b^{LR, \tilde{\chi}^+} = \frac{h_b h_t}{16\pi^2} \mu (T_1^{FC} + T_2^{FV} + T_3^{FC} + T_3^{FV} + \dots), \quad (24)$$

where the explicit result for the terms T_i^x (see Ref. [1]) is lengthy and therefore not shown here.

5. Numerical analysis

To demonstrate the effects of QFV we have chosen a reference scenario with strong $\tilde{c} - \tilde{t}$ mixing in both $h^0 \rightarrow b\bar{b}$ and $h^0 \rightarrow c\bar{c}$ decays. The corresponding MSSM parameters at $Q = 1$ TeV are shown in Table 1. This scenario satisfies all present experimental and theoretical constraints explicitly listed in Ref. [1]. The resulting physical masses of the particles are shown in Table 2. The flavour decomposition of the up-type squarks \tilde{u}_i , $i = 1, \dots, 6$ is shown in Table 3. In the following, unless specified otherwise, we show various parameter dependences of the relative to the SM width $\Gamma/\Gamma^{\text{SM}} - 1$ for $\Gamma(h^0 \rightarrow b\bar{b})$ and $\Gamma(h^0 \rightarrow c\bar{c})$ with all other parameters fixed as in Table 1.

In Fig. 2 the dependence on the QFV parameters δ_{23}^{uRL} and δ_{23}^{uLR} is shown. It is seen that in the case of $b\bar{b}$

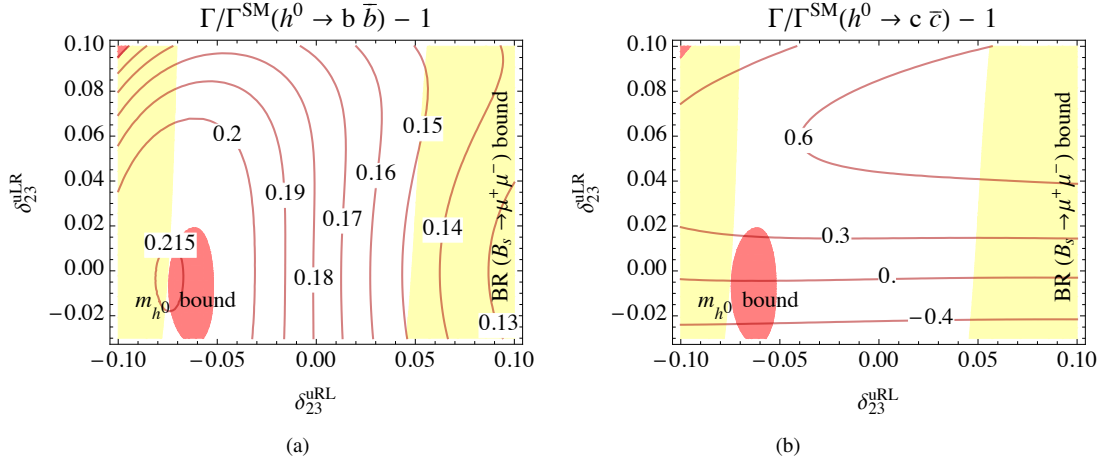


Figure 2: Contours of the deviation (a) $\Gamma/\Gamma^{\text{SM}}(h^0 \rightarrow b\bar{b}) - 1$ and (b) $\Gamma/\Gamma^{\text{SM}}(h^0 \rightarrow c\bar{c}) - 1$ in the δ_{23}^{uRL} - δ_{23}^{uLR} plane for $\delta_{23}^{uRR} = 0.5$ and $\delta_{23}^{LL} = 0$.

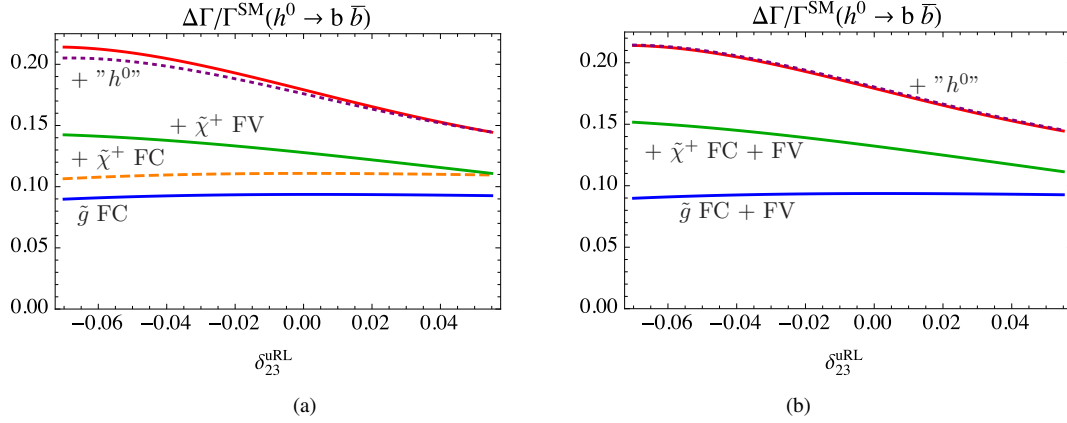


Figure 3: Contours of the deviation (a) $\Gamma/\Gamma^{\text{SM}}(h^0 \rightarrow b\bar{b}) - 1$ in the δ_{23}^{uRR} - δ_{23}^{uRL} plane for $\delta_{23}^{LL} = 0$ and $\delta_{23}^{uLR} = 0$ and (b) $\Gamma/\Gamma^{\text{SM}}(h^0 \rightarrow c\bar{c}) - 1$ in the δ_{23}^{uRR} - δ_{23}^{uLR} plane for $\delta_{23}^{LL} = 0$ and $\delta_{23}^{uRL} = 0.02$.

(Fig. 2(a)) the variation due to correlated $\tilde{c}_R - \tilde{t}_L$ and $\tilde{c}_L - \tilde{t}_R$ mixing can vary up to $\sim 7\%$ in the region allowed by the constraints. Comparing Fig. 2(a) with Fig. 2(b) one can see that there exist regions where both widths considered simultaneously deviate significantly from their SM prediction. Hence $\Gamma(h^0 \rightarrow b\bar{b})$ tends to depend more on $\tilde{c}_R - \tilde{t}_L$ mixing, while $\Gamma(h^0 \rightarrow c\bar{c})$ depends more on $\tilde{c}_L - \tilde{t}_R$ mixing.

In Section 4.1, in agreement with our results in Ref. [3], we have shown that in the case of $c\bar{c}$ the deviation from the SM is entirely due to QFV. However, it is known that in the MSSM $\Gamma(h^0 \rightarrow b\bar{b})$ can differ considerably from the SM due to quark-flavour conserving (QFC) contributions [5]. In Fig.3(a) the individual one-loop contributions to $\Gamma(h^0 \rightarrow b\bar{b})$ as a func-

tion of δ_{23}^{uRL} are shown. The top curve shows the full one-loop contribution to the width with no approximation. It is seen that the main one-loop contribution to $\Gamma(h^0 \rightarrow b\bar{b})$ comes from QFC gluino and chargino exchange. Nevertheless, there exist a region for large and negative δ_{23}^{uRL} where the QFV component can be comparable with the QFC component. The QFV component is mainly due to chargino exchange which involves mixing in the \tilde{u} -sector. On the other hand, the QFV gluino exchange, which plays a major role in the $c\bar{c}$ case, in the $b\bar{b}$ case involves \tilde{d} quarks whose mixing is strongly suppressed, and hence, this contribution is very small and thus not considered here. It is also interesting that the QFV component receives a large contribution from the dependence of $\Gamma^{g,\text{impr}}$ on the Higgs mass and the angle

Table 1: Reference scenario: shown are the basic MSSM parameters at $Q = 1$ TeV, except for m_{A^0} which is the pole mass of A^0 , with $T_{U33} = 1450$ GeV (corresponding to $\delta_{33}^{uRL} = 0.1$). All other squark parameters not shown here are zero.

M_1	M_2	M_3
400 GeV	800 GeV	2000 GeV

μ	$\tan\beta$	m_{A^0}
500 GeV	30	1500 GeV

	$\alpha = 1$	$\alpha = 2$	$\alpha = 3$
$M_{Q\alpha\alpha}^2$	3200^2 GeV^2	1550^2 GeV^2	1100^2 GeV^2
$M_{U\alpha\alpha}^2$	3200^2 GeV^2	2800^2 GeV^2	2050^2 GeV^2
$M_{D\alpha\alpha}^2$	3200^2 GeV^2	3000^2 GeV^2	2500^2 GeV^2

δ_{23}^{LL}	δ_{23}^{uRR}	δ_{23}^{uRL}	δ_{23}^{uLR}
0	0.8	0.02	0.02

Table 2: Physical masses in GeV of the particles for the scenario of Table 1.

$m_{\tilde{\chi}_1^0}$	$m_{\tilde{\chi}_2^0}$	$m_{\tilde{\chi}_3^0}$	$m_{\tilde{\chi}_4^0}$	$m_{\tilde{\chi}_1^+}$	$m_{\tilde{\chi}_2^+}$
395	507	511	845	501	845

m_{h^0}	m_{H^0}	m_{A^0}	m_{H^\pm}
125	1500	1500	1503

$m_{\tilde{g}}$	$m_{\tilde{u}_1}$	$m_{\tilde{u}_2}$	$m_{\tilde{u}_3}$	$m_{\tilde{u}_4}$	$m_{\tilde{u}_5}$	$m_{\tilde{u}_6}$
2103	996	1176	1578	3214	3217	3327

$m_{\tilde{d}_1}$	$m_{\tilde{d}_2}$	$m_{\tilde{d}_3}$	$m_{\tilde{d}_4}$	$m_{\tilde{d}_5}$	$m_{\tilde{d}_6}$
1128	1579	2515	3012	3211	3218

Table 3: Flavour decomposition of \tilde{u}_i , $i = 1, \dots, 6$ for the scenario of Table 1. Shown are the squared coefficients.

	\tilde{u}_L	\tilde{c}_L	\tilde{t}_L	\tilde{u}_R	\tilde{c}_R	\tilde{t}_R
\tilde{u}_1	0	0.002	0.25	0	0.228	0.52
\tilde{u}_2	0	0	0.749	0	0.086	0.165
\tilde{u}_3	0.051	0.946	0.001	0	0	0
\tilde{u}_4	0.95	0.05	0	0	0	0
\tilde{u}_5	0	0	0	1	0	0
\tilde{u}_6	0	0	0	0	0.69	0.31

α , which depend on the QFV parameters. In Figs. 3(a) and 3(b) this contribution is denoted by " h^0 ". Note that m_{h^0} as well as $\sin\alpha$ already appear in the kinematics factor at tree level, see eq. (8). The total QFV contribution to $\Gamma/\Gamma^{\text{SM}}(h^0 \rightarrow b\bar{b})$ can be as large as $\sim 8\%$ at a certain point. Fig.3(b), where no MI is used, demonstrates in addition the quality of the total approximated result obtained in [1], see eq. (4.38) therein. A numerical comparison of the different MI orders shows that the MI formulas converge fast for \tilde{g} FC and $\tilde{\chi}^+$ FC, but not for $\tilde{\chi}^+$ FV, compare Fig. 3(a) with Fig. 3(b). The difference between the dotted curve and the upper curve in Fig. 3(a), therefore, is mainly due to the relatively slow MI convergence of the $\tilde{\chi}^+$ FV contribution.

Although the decay $h^0 \rightarrow b\bar{b}$ is dominant, the measurement of its branching ratio and width at the LHC will be hard due to the huge QCD background. In any case, high luminosity at LHC would be needed [6]. A model independent and precise measurement of $B(h^0 \rightarrow b\bar{b})$ and $\Gamma(h^0 \rightarrow b\bar{b})$ would be possible at a e^+e^- linear collider such as ILC [7].

6. Conclusions

We have studied the decays $h^0 \rightarrow b\bar{b}$ and $h^0 \rightarrow c\bar{c}$ at full one-loop level in the MSSM with quark-flavour mixing in the heavy squark sector. The dominant contributions with gluino and chargino exchange are calculated in the mass insertion approximation, using the Flavour Expansion Theorem. Both widths, $\Gamma(h^0 \rightarrow b\bar{b})$ and $\Gamma(h^0 \rightarrow c\bar{c})$, can deviate from the SM significantly within the allowed parameter region. In the $c\bar{c}$ case the deviation is mainly due to the MSSM QFV parameters. In the $b\bar{b}$ case the deviation is mainly due to the MSSM QFC parameters, but nevertheless at certain parameter regions the the QFV parameters can cause fluctuations of $\Gamma(h^0 \rightarrow b\bar{b})$ up to $\sim 7\%$, with similar large contribution coming from the Higgs parameters dependence.

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